A Realization of Effective SUSY with Strong Unification

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Abstract

A natural model of realizing the effective supersymmetry is presented. Two sets of the Standard Model-like gauge group $G_1 \times G_2$ are introduced, where $G_i = SU(3)_i \times SU(2)_i \times U(1)_i$, which break diagonally to the Standard Model gauge group at the energy scale $M \sim 10^7$ GeV. Gauge couplings in G_1 are assumed much larger than that in G_2. Gauge mediated supersymmetry breaking is adopted. The first two generations (third one) are charged only under G_1 (G_2). The effective supersymmetry spectrum is obtained. How to reproduce realistic Yukawa couplings is studied. Fine-tuning for an 126 GeV Higgs is much reduced by the large A term due to direct Higgs-messenger interaction. Finally, G_2 is found to be a non-trivial realization of the strong unification scenario in which case we can predict $\alpha_s(M_Z)$ without real unification

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I. INTRODUCTION

A Standard Model (SM)-like Higgs particle of 126 GeV has been discovered at the LHC [1]. If we are insisting on naturalness of the SM, this discovery strengthens motivation for the low energy supersymmetry (SUSY) which stabilizes the Higgs mass at the electroweak (EW) scale. However, there is yet no definite sign of sparticles after the integrated luminosity has reached 20 fb^{-1} at $\sqrt{s} = 8$ TeV. This null result for the SUSY search sets a lower bound for the first two generation squarks – $m_{\tilde{Q}_{1,2}} > 1$ TeV, while stops/sbottoms with a mass about 500 GeV are still allowed [2]. Noticing that naturalness sets an upper bound for sparticle masses as 1 TeV.

SUSY, if it is relevant to the EW physics, should be beyond its simplest version. Actually it was noted long time ago that naturalness only requires the third generation sfermions and particles (gauginos and Higgsinos) that interact significantly with the Higgs to have sub-TeV masses, while the first two generation sparticles can be heavy up to 20 TeV [3–5]. Such sparticle spectra also alleviate the SUSY FCNC problem. This phenomenological scenario is dubbed Effective SUSY by Cohen *et al.* [5]. Nowadays this Effective SUSY has become one of the main scenarios to reconcile naturalness with the null SUSY search [6].

In this paper, we realize Effective SUSY through modifying models of Refs.[7, 8]. In those models we have introduced two sets of SM-like gauge groups $G_1 \times G_2$ where $G_i = SU(3)_i \times SU(2)_i \times U(1)_i$. At the TeV scale, G_1 is strongly and G_2 is weakly interacting, respectively. They break diagonally to the SM gauge group. SUSY breaking is due to $G_1 \times G_2$ gauge mediation. Furthermore, G_2 is of strong unification [8], namely its gauge coupling constants have a common Landau pole at the unification scale [9–11]. In that model, all the three generations were put in G_2, that did not result in any sparticle splitting. To make sparticle splitting, the first two generations have to be treated differently from the third one. In this work, the first two generations are put in G_1 and the third in G_2.

Among other things, we need to solve the following problems. First, to generate Yukawa interactions between the first two generations and the third generation. This is because three generations are in different gauge sectors from the beginning. Then, to reduce fine tuning of the 126 GeV Higgs. In conventional gauge mediated SUSY breaking (GMSB) [12], fine-tuning seemed unavoidable for a 126 GeV Higgs. Furthermore, to re-examine strong unification. This is needed due to all the changes in the particle content, and conditions of

strong unification are quite subtle.

We notice that a similar model was proposed by Craig *et al.* [13, 14] in which the first two generations were also put in G_1 and the third one in G_2. However, there are several main differences. First is about G_1. In the model of [14], G_1 gauge interactions are weakly coupled at the TeV scale, whereas our G_1 is superstrong. Second is for G_2. Their G_2 gauge interactions unify weakly in the sense of ordinary grand unification, our G_2 is of strong unification [9–11]. Third is about GMSB. They only use a messenger for G_1, and we have messengers both for G_1 and G_2. And finally we need to use direct Higgs-messenger interaction to reduce fine-tuning of a 126 GeV Higgs. These differences make this model qualitatively different.

There were a few other ways to realize Effective SUSY [15, 16]. Some of them used an extra $U(1)$ gauge group which contributes larger masses to the first two generation sparticles than to the third generation ones. Usually, this $U(1)$ symmetry suppresses Yukawa couplings for the first two generations compared to that for the third generation, giving an explanation for the fermion mass hierarchy. Some other works assume particular boundary conditions at a high scale, then employ the technique of renormalization group [16].

The paper is organized as follows. Our model will be given in the next section, where flavor physics for fermions and the problem of naturalness in light of a 126 GeV Higgs will be discussed. After all particle contents and the mass spectrum have been fixed, the prediction for $\alpha_s(M_Z)$ will be calculated by means of strong unification in Section III. The final section summarizes our results and gives discussions.

II. THE MODEL

We consider a SUSY model with two sets of the SM-like gauge group G_1 and G_2 where $G_i = SU(3)_i \times SU(2)_i \times U(1)_i$. The first two and the third generation of matter transform under G_1 and G_2, respectively. The two Higgs doublets H_u and H_d are in G_2. The other fields include SUSY breaking messengers and the Higgs fields which break $G_1 \times G_2$ into the SM. For convenience, we will use field representations under $SU(5)$ to illustrate their representations under $SU(3) \times SU(2) \times U(1)$.

The GMSB mechanism is employed. Two sets of messenger fields T_1 ($\bar{T}_1$) and T_2 ($\bar{T}_2$) are introduced. They transform nontrivially under G_1 and G_2, respectively. Without losing

generality, we will focus on the quark/squark sector. At the scale of SUSY breaking, squarks have following masses,

$$m^2_{\tilde{Q}_{1,2}} \sim (\frac{g_1^2}{16\pi^2}\frac{F}{M})^2 , \quad \text{and} \quad m^2_{\tilde{Q}_3} \simeq (\frac{g_2^2}{16\pi^2}\frac{F}{M})^2 . \tag{1}$$

M stands for the messenger scale and $\sqrt{F}$ is the measure of SUSY breaking. g_1 and g_2 represent coupling constants for G_1 and G_2 respectively. Therefore, we can realize Effective SUSY sparticle spectrum by requiring g_1 to be much larger than g_2. Note that Eq.(1) for $m^2_{\tilde{Q}_{1,2}}$ is not exact, because g_1 is too large.

A pair of Higgs fields Φ and $\bar{\Phi}$ charged under $G_1 \times G_2$ as $\mathbf{5} \times \bar{\mathbf{5}}$ and $\bar{\mathbf{5}} \times \mathbf{5}$ is introduced. Φ and $\bar{\Phi}$ have a mass M_Φ, and they obtain vacuum expectation values (VEVs) as $\langle \Phi \rangle = \langle \bar{\Phi} \rangle = V I_2 \times I_3$, where $V \sim M_\Phi$, I_2 and I_3 are the unit matrix in the subspace of $SU(2)_1 \times SU(2)_2$ and $SU(3)_1 \times SU(3)_2$, respectively. As a result, $G_1 \times G_2$ break diagonally to the SM gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y$. Below the scale of M_Φ, the effective theory of this model looks like Effective SUSY with following relations among gauge coupling constants,

$$\frac{1}{g_s^2} = \frac{1}{g_{s1}^2} + \frac{1}{g_{s2}^2} , \quad \frac{1}{g^2} = \frac{1}{g_1^2} + \frac{1}{g_2^2} , \quad \frac{1}{g'^2} = \frac{1}{g_1'^2} + \frac{1}{g_2'^2} . \tag{2}$$

Because of the relation $g_1^2 \gg g_2^2$, SM gauge couplings are almost fully determined by that in G_2. More details about breaking of $G_1 \times G_2$ can be found in Ref. [8]. While this model in many aspects is similar to that of Ref. [8], new features come in because of separation of three generations.

Before $G_1 \times G_2$ breaking, three generations are put in different gauge sectors, so there is no marginal operators giving Yukawa couplings between the first two generation matter and the third generation ones. However, higher-dimensional operators such as $H_d Q_3 \Phi_{d\bar{d}} \bar{d}_{1(2)}$ are allowed by $G_1 \times G_2$ symmetry. Here $\Phi_{d\bar{d}}$ is the component of the Higgs Φ with quantum numbers $(3, 1, -\frac{1}{3}) * (\bar{3}, 1, \frac{1}{3})$ under $SU(3)_1 \times SU(2)_1 \times U(1)_1 \times SU(3)_2 \times SU(2)_2 \times U(1)_2$. This kind of operators can be produced by integrating out appropriate heavy fields, just like the Froggatt-Nielsen mechanism [17]. As a result, they are suppressed by M (mass scale for the heavy fields that have been integrated out) $\frac{1}{M} H_d Q_3 \Phi_{d\bar{d}} \bar{d}_{1(2)}$. When $G_1 \times G_2$ is spontaneously broken i.e. $\Phi_{d\bar{d}}$ gets VEV, there will be terms like $\frac{V}{M} H_d Q_3 \bar{d}_{1(2)}$ that lead to Yukawa interactions between 1st/2nd generation fermions and 3rd generation fermions. Taking $\frac{V}{M}$ a small quantity $\sim 0.1-0.01$, the mass hierarchy between the third generation and the first two is obvious. Roughly speaking, this paves the way to obtain the realistic fermion

mass pattern, mixing and CP violation. We can say hierarchy among three generation fermions and that among three generation sfermions are closely connected to each other in this model.

This approach has been discussed in Ref. [14]. There, a vector-like **5** representation is introduced as the mediator that is integrated out. Similarly, in G_2, we introduce a full vector-like generation $(L, \bar{d})$ and $(Q, \bar{u}, \bar{e})$ as representation $\bar{\mathbf{5}}$ and $\mathbf{10}$, respectively. Masses of these vectorlike fields are taken to be of the same order M_Ψ.

In MSSM, the SM-like Higgs has the following mass after including one-loop radiative corrections induced by stops,

$$m_h^2 = m_Z^2 \cos^2 2\beta + \frac{3m_t^4}{4\pi^2 v^2}[\log \frac{m_{\tilde{t}}^2}{m_t^2} + \frac{X_t^2}{m_{\tilde{t}}^2}(1 - \frac{X_t^2}{12 m_{\tilde{t}}^2})], \tag{3}$$

where $X_t = A_t - \mu \cot \beta$. To obtain a 126 GeV Higgs, we either need multi-TeV stops which lead to severe fine-tuning, or turn to the large A_t scenario in which case sub-TeV stops are enough. In conventional GMSB, contribution of A_t is negligible and fine-tuning induced by too heavy stops seems to be unavoidable. To preserve naturalness, we will produce a large A_t term by extending conventional GMSB.

We choose the messenger $T_2(\bar{T}_2)$ to be the representation $\mathbf{10}$ $(\bar{\mathbf{10}})$. It is found that [18], a large A_t term can be produced without a large Higgs mass. This is due to direct interaction between the Higgs and the messenger in the superpotential,

$$y H_u T_2^Q T_2^{\bar{u}}\,, \tag{4}$$

where T_2^Q and $T_2^{\bar{u}}$ are components of T_2 that have same gauge quantum numbers as Q_3 and $\bar{u}_3$, respectively. It should be pointed out that this term is the only one of direct interaction between messengers and ordinary matters by employing the messenger parity. One-loop contribution to the soft term $H_u \tilde{Q}_3 \tilde{\bar{u}}_3$ is extracted from wave function renormalization for the superfield H_u [19],

$$A_t \sim -\frac{y^2 y_t}{16\pi^2} \frac{F}{M}\,. \tag{5}$$

In Ref. [20], the same way was used to produce a large A_t term except that there it was Q_3 and $\bar{u}_3$ instead of H_u that interact directly with the messenger. More issues about this way of producing large A_t term can be found in [18][21].

III. STRONG UNIFICATION

We will consider unification of gauge coupling constants. Because the SM gauge couplings are almost fully determined by those of G_2, what we will really care about, above $G_1 * G_2$ breaking scale, is the G_2 gauge coupling constants. So far quite a few new fields have been introduced. The particle content and the mass spectrum are summarized in the following with emphasis on that which are charged under G_2. There is only one chiral generation in G_2. The two Higgs doublets are in G_2. The bi-fundamental Higgs Φ ($\bar{\Phi}$) is charged under $SU(5)_1 \times SU(5)_2$ as $\mathbf{5} \times \bar{\mathbf{5}}$ ($\bar{\mathbf{5}} \times \mathbf{5}$) with a mass M_Φ. The messengers T_2 ($\bar{T}_2$) is charged under G_2 as representation $\mathbf{10}$ ($\bar{\mathbf{10}}$) with a mass M. Besides, there is an extra vector-like generation charged under G_2 as representation $\bar{\mathbf{5}} + \mathbf{10}$ ($\mathbf{5} + \bar{\mathbf{10}}$) with a mass M_Ψ. For simplicity and definiteness, we will identify M_Φ and M_Ψ with M in the following analysis.

Below $G_1 * G_2$ breaking scale M, this model is just that of the minimal SUSY SM, so SM gauge couplings running can be calculated in the usual way. At the scale M, SM gauge couplings are identified with that in G_2. Above the scale M, since there are so many complete representations of $SU(5)_2$, while the unification energy scale does not change, gauge couplings in G_2 grow so fast that they may come across their Landau poles as they evolve to the unification scale. The situation, where gauge couplings reach their common Landau pole, is named as strong unification [8–10]. Using strong unification to predict gauge couplings at the EW scale seems unreasonable because of the strong coupling domain where the perturbative method is not reliable. However, as shown in Ref. [10], ratios of gauge couplings in G_2 will reach their infra-fixed points at the scale M. Thus, we can determine SM gauge couplings at the EW scale, with ratios of gauge couplings in G_2 at the scale M as boundary condition where perturbative calculation already works.

Boundary conditions for gauge couplings of G_2 at the scale M satisfy the following relation [10],

$$\alpha_2'(M) b_2' = \alpha_2(M) b_2 = \alpha_{s2}(M) b_{s2}, \tag{6}$$

where $\alpha_2'(M) = \dfrac{g_2'^2(M)}{4\pi}$ and so forth. $b_2' = \dfrac{73}{5}$, $b_2 = 9$ and $b_{s2} = 5$ are one-loop beta functions above the scale M for g_2', g_2 and g_{s2}, respectively. In the following, we will first calculate the prediction for $\alpha_s(M_Z)$ in a simply way to illustrate the usage of strong unification, and then take into consideration two-loop contributions and low-scale threshold effects induced by sparticles' mass splitting among different generations.

In the first case, $\alpha'(M)$ and $\alpha(M)$ can be determined by $\alpha'(M_Z)$ and $\alpha(M_Z)$ through the following equations,

$$\alpha'^{-1}(M) = \alpha'^{-1}(M_Z) - \frac{b'}{2\pi}\ln\frac{M}{M_Z} \tag{7}$$

$$\alpha^{-1}(M) = \alpha^{-1}(M_Z) - \frac{b}{2\pi}\ln\frac{M}{M_Z}. \tag{8}$$

With Eq.(6), we can get $M \sim 10^8$ GeV and $\alpha_s^{-1}(M) = 15.17$. Finally, $\alpha_s(M_Z)$ is calculated to be 0.119 as follows,

$$\alpha_s^{-1}(M_Z) = \alpha_s^{-1}(M) + \frac{b_s}{2\pi}\ln\frac{M}{M_Z}. \tag{9}$$

After inclusion of low scale threshold effects and dominated two-loop contributions, Eqs.(7-9) will be replaced by the following equations,

$$\begin{aligned}\alpha'^{-1}(M) = {}& \alpha'^{-1}(M_Z) - \frac{\tilde{\tilde{b}}'}{2\pi}\ln\frac{m_{\tilde{Q}_3}}{M_Z} - \frac{\tilde{b}'}{2\pi}\ln\frac{m_{\tilde{Q}_{1,2}}}{m_{\tilde{Q}_3}} - \frac{b'}{2\pi}\ln\frac{M}{m_{\tilde{Q}_{1,2}}}\\ &-\frac{1}{4\pi}\frac{b_{11}}{b'}\ln\frac{\alpha'(M)}{\alpha'(M_Z)} - \frac{1}{4\pi}\frac{b_{12}}{b}\ln\frac{\alpha(M)}{\alpha(M_Z)} - \frac{1}{4\pi}\frac{b_{13}}{b_s}\ln\frac{\alpha_s(M)}{\alpha_s(M_Z)},\end{aligned} \tag{10}$$

$$\begin{aligned}\alpha^{-1}(M) = {}& \alpha^{-1}(M_Z) - \frac{\tilde{\tilde{b}}}{2\pi}\ln\frac{m_{\tilde{Q}_3}}{M_Z} - \frac{\tilde{b}}{2\pi}\ln\frac{m_{\tilde{Q}_{1,2}}}{m_{\tilde{Q}_3}} - \frac{b}{2\pi}\ln\frac{M}{m_{\tilde{Q}_{1,2}}}\\ &-\frac{1}{4\pi}\frac{b_{21}}{b'}\ln\frac{\alpha'(M)}{\alpha'(M_Z)} - \frac{1}{4\pi}\frac{b_{22}}{b}\ln\frac{\alpha(M)}{\alpha(M_Z)} - \frac{1}{4\pi}\frac{b_{23}}{b_s}\ln\frac{\alpha_s(M)}{\alpha_s(M_Z)},\end{aligned} \tag{11}$$

$$\begin{aligned}\alpha_s^{-1}(M) = {}& \alpha_s^{-1}(M_Z) - \frac{\tilde{\tilde{b}}_s}{2\pi}\ln\frac{m_{\tilde{Q}_3}}{M_Z} - \frac{\tilde{b}_s}{2\pi}\ln\frac{m_{\tilde{Q}_{1,2}}}{m_{\tilde{Q}_3}} - \frac{b_s}{2\pi}\ln\frac{M}{m_{\tilde{Q}_{1,2}}}\\ &-\frac{1}{4\pi}\frac{b_{31}}{b'}\ln\frac{\alpha'(M)}{\alpha'(M_Z)} - \frac{1}{4\pi}\frac{b_{32}}{b}\ln\frac{\alpha(M)}{\alpha(M_Z)} - \frac{1}{4\pi}\frac{b_{33}}{b_s}\ln\frac{\alpha_s(M)}{\alpha_s(M_Z)},\end{aligned} \tag{12}$$

with

$$\begin{array}{llllll} \tilde{\tilde{b}}' = \frac{41}{10} & \tilde{b}' = \frac{79}{15} & b' = \frac{33}{5} & b_{11} = \frac{199}{25} & b_{12} = \frac{27}{5} & b_{13} = \frac{88}{5} \\ \tilde{\tilde{b}} = -\frac{19}{6} & \tilde{b} = -\frac{1}{3} & b = 1 & b_{21} = \frac{9}{5} & b_{22} = 25 & b_{23} = 24 \\ \tilde{\tilde{b}}_s = -7 & \tilde{b}_s = -\frac{13}{3} & b_s = -3 & b_{31} = \frac{11}{5} & b_{32} = 9 & b_{33} = 14 \end{array} \tag{13}$$

It is found that $\alpha_s(M_Z) \sim 0.117$ and $M \sim 10^7$ GeV, when $m_{\tilde{Q}_3}$ and $m_{\tilde{Q}_{1,2}}$ take typical value 1 TeV and 10 TeV respectively. This value is very close to world average value 0.1184 ± 0.0007 [22].

In the above discussion, we have taken the limit $\frac{g_1^2}{g_2^2}$ goes to infinity so that SM gauge coupling can be identified with that in G_2 at the Higgsing scale. There will be several

percents uncertainty in this identification, if we take into consideration that $\frac{g_1^2}{g_2^2}$ is finite which is several tens. This uncertainty will affect the prediction for $\alpha_s(M_Z)$ substantially. For example, with a typical value $20 \sim 40$ for $\frac{g_1^2}{g_2^2}$, $\alpha_s(M_Z)$ will have an uncertainty about 0.005. However, if three gauge couplings in G_1 sector has the same ratio with the counterparts in G_1 sector,

$$\frac{g_1'}{g_2'} = \frac{g_1}{g_2} = \frac{g_{s1}}{g_{s2}} \tag{14}$$

this uncertainty will disappear. This is because the boundary condition Eq.(6) just depends on the ratio of gauge coupling constants.

IV. SUMMARY AND DISCUSSION

In summary, we have presented a model of realizing the Effective SUSY. Two sets of the SM-like gauge group $G_1 \times G_2 = SU(3)_1 \times SU(2)_1 \times U(1)_1 \times SU(3)_2 \times SU(2)_2 \times U(1)_2$ have been introduced which break diagonally to the SM gauge group at the energy scale $M \sim 10^7$ GeV. Gauge couplings in G_1 have been assumed much larger than that in G_2. GMSB has been adopted. The first two generations (third one) are charged only under G_1 (G_2). The Effective SUSY spectrum has been obtained naturally. Fine-tuning for an 126 GeV Higgs is much reduced. With all the fields necessary and their masses fixed, $\alpha_s(M_Z)$ can be predicted in the scenario of strong unification.

Compared to our previous works [7, 8], in addition to the Effective SUSY spectrum, following new features have arisen.

(1) An extra vector-like generation charged under G_2 has been introduced as mediator, so as to reproduce realistic Yukawa couplings between the first two generations and the third generation, i.e. the suitable fermion mass hierarchy and the CKM mixing matrix.

(2) Fine-tuning for an 126 GeV Higgs is much reduced by a large A_t term produced by direct Higgs-messenger interaction, because the messenger for G_2 has been specified to be a **10** representation under the $SU(5)$, which is absent in conventional GMSB.

The following three main aspects clarify differences of Ref. [14] from our model.

(a) In Ref. [14], gauge couplings in G_1 and G_2 were comparable. Only a messenger for G_1 was introduced, and the third generation sparticles could feel SUSY breaking only after the breaking of $G_1 \times G_2$, so that $m_{\tilde{Q}_3}$ was suppressed by an additional factor $\frac{V}{M}$ in comparison

with $m_{\tilde{Q}_{1,2}}$.

(b) There was no need to produce a large A_t term in Ref. [14]. Due to the comparability of gauge couplings in G_1 and G_2, and the low scale of $M_\Phi \sim 10^4$ GeV, non-decoupling D term contribution to the Higgs mass could be significant.

(c) There, unification of gauge couplings was "weak" in the sense of in comparison with strong unification.

Here comes our final remarks. First, it is worth pointing out that despite the term strong "unification", $SU(3) \times SU(2) \times U(1)$ do not necessarily unify into a larger simple group, so that there can be no proton decay at all, and thus there is no the so-called doublet-triplet splitting problem. Because $g_2 \ll g_1$, the G_2 unification scale is the same as that of the traditional GUT, namely about 3×10^{16} GeV. Second, gauge couplings in G_1 are expected to be a realization of GUT. We have not studied that much because it does not affect our physical results on one hand and the couplings are too strong to use perturbation method on the other hand. Third, LHC has set a lower bound on the gluino mass $m_{\tilde{g}} > 1$ TeV [2], and this bound would also apply to $m_{\tilde{Q}_3}$ in traditional GMSB. This is not the case for this model, because $G_1 \times G_2$ breaking also contributes the gluino an additional mass $\sim \dfrac{g_2^2 V^2}{M}$ from mixing with the fermionic component of Φ. This contribution is expected to be larger than the purely soft mass $\dfrac{g_2^2}{16\pi^2}\dfrac{F}{M}$ [8]. Besides, Higgs-mediated SUSY breaking contribution reduces $m_{\tilde{Q}_3}$. In a word, this model allows an interesting mass pattern $m_{\tilde{Q}_{1,2}} \gg m_{\tilde{g}} \gg m_{\tilde{Q}_3}$. Finally, in this work we have taken that the first two generations in the same gauge group. It is imaginable that we can introduce one more version of the SM group to split these two generations further. Namely we may expect a model of $[SU(3) \times SU(2) \times U(1)]^3$, which first breaks into $G_1 \times G_2$ at some higher energy scale.

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